

Nucleon polarization in deuteron electrodisintegration

B. Mosconi

Dipartimento di Fisica, Università di Firenze, I-50125 Firenze, Italy
Istituto Nazionale di Fisica Nucleare, Sezione di Firenze, I-50125 Firenze, Italy

J. Pauschenwein

Institut für Theoretische Physik, Universität Graz, A-8010 Graz, Austria

P. Ricci

Istituto Nazionale di Fisica Nucleare, Sezione di Firenze, I-50125 Firenze, Italy

Outgoing nucleon polarization in exclusive deuteron electrodisintegration is studied at the quasi-elastic peak in the standard theory with emphasis on the effect of nucleonic and pionic relativistic corrections. The cases of polarized beam or/and polarized target are considered. Sizeable relativistic effects are pointed out in several polarization components. The sensitivity of nucleon polarization to the neutron charge form factor G_E^n is discussed. In particular, it is shown that the longitudinal component of neutron polarization with vector polarized deuterons is as sensitive to G_E^n as the sideways beam polarization transfer.

I. INTRODUCTION

In this paper we shall study the outgoing nucleon polarization in exclusive deuteron electrodisintegration with beam and target polarized, taking into account the relativistic corrections (RC) to the standard theory in a perturbation expansion approach [1-6]. Such theory, which in addition to the non relativistic impulse approximation (IA) considers the effects of meson exchange (MEC) and isobar excitation (IC) currents, has proved to be rather successful for the electromagnetic (EM) interactions in few-body nuclei at low and medium energy-momentum transfers [7]. However, consideration of nucleonic and pionic RC turns out essential for explaining some observables in elastic electron scattering as well as in photo- and electro-disintegration. A well known case concerns deuteron photodisintegration whose differential cross section is heavily affected by RC particularly in the forward direction, as our group first pointed out [8] to resolve the long standing discrepancy in the energy spectrum of the 0° cross section measured at Mainz in 1976 [9]. Obviously, this large relativistic effect in forward and backward direction strongly modifies the shape of the differential cross section [8 b]. Indeed, the recent data on the proton angular distribution taken with quasi monochromatic photon beams or with absolutely calibrated neutron beams in the time reversed reaction, are in good agreement with the relativistic predictions [10]. For an exhaustive survey of deuteron photodisintegration we refer to the review paper by Arenhövel and Sanzone [11].

A second and more recent example is that of the longitudinal-transverse structure function f_{LT} and thus, of the ϕ -asymmetry (A_ϕ) of the $d(e, e'p)n$ cross section. A sizeable relativistic effect in f_{LT} has been found by two of us [12] in a paper devoted to the influence of RC on the exclusive cross section and has been recently verified by van der Schaar et al. [13] in a three-fold separation experiment at NIKHEF-K. We would like to note that our results of A_ϕ [14] are hardly distinguishable from those obtained by Hummel and Tjon [15], as reported in Ref.[13], in a complete covariant model based on the Bethe-Salpeter equation. This makes us confident in the reliability of our perturbative calculations limited to terms of order $(v/c)^2$ in the energy and momentum transfer region of the present accelerators. Of course, the situation may well change at higher energies where higher order terms become necessary, because the perturbative expansion in power of (p/m) seems to converge badly [16].

An even stronger relativistic effect on f_{LT}^h , the fifth structure function arising for polarized electrons, which was predicted in Ref.[12], remains to be experimentally confirmed. Note that an experiment on f_{LT}^h or equivalently of the electron asymmetry of the cross section requires an out-of-plane spectrometer in addition to a polarized

electron beam.

Coming back to the nucleon polarization , its first measurement in exclusive deuteron break-up has been carried out very recently using the longitudinally polarized electron beam of the MIT - Bates Lab. [17] and a liquid deuterium target, but the results are still to be published. Such experiment, prohibitive in the past because of the low duty cycle of the previous generation of accelerators in spite of the improvements in nucleon polarimeters, has been made possible by the upgrading of the accelerator and by the advances in the polarized electron source technology. Note that there has not yet been a measurement of \mathbf{P}_0 , the polarization with unpolarized beam and target. The reason for choosing the beam polarization transfer \mathbf{P}_0^h over \mathbf{P}_0 in the first measurement of neutron polarization lies in the interest of the sideways component of \mathbf{P}_0^h for emitted neutron ($P_{0x}^h(n)$) as a source of information about the poorly known charge form factor of the neutron (G_E^n) as first suggested by Arnold et al. [18] and studied by several authors [19-22]. In particular, Arenhövel and collaborators [21,22] have shown that $P_{0x}^h(n)$ is almost independent of NN potential models and of meson exchange effects in the parallel kinematics. Unlike previous papers on outgoing nucleon polarization [18-24] we shall consider in detail also the target polarization transfer, i.e. the polarization induced by polarized deuterons, because the rapid progress in the production of polarized deuteron targets makes it possible to foresee their measurement in the near future. For completeness, we shall also consider the beam-target polarization transfer.

Much effort for electron scattering from polarized targets is concentrated in the internal gas target experiments [25] which have the advantages of purity of the targets and of rapid polarization reversal while an acceptable luminosity is obtained even with thin polarized gas targets thanks to the high current of the beam. Presently, the devices used for feeding polarized deuteron targets are the atomic beam sources which allow to reach vector and tensor polarization close to unity. Unfortunately, several depolarizing effects (cell wall bounces, electron beam magnetic field,...) may considerably reduce the polarization degree in the storage cell. This technique has been successfully employed at Novosibirsk VEPP-3 storage ring for an experiment on tensor analysing power in elastic electron deuteron scattering [26] and also for the first measurement of the tensor asymmetries of the exclusive inelastic cross section as reported in Ref.[27]. The major problem with the atomic beam sources is the low intensity of polarized atoms injected into the storage cell. More than an order of magnitude can be gained with the laser driven sources when operating in a high magnetic field so compensating the figure of merit for the lower values of deuteron polarization achievable [28]. In conclusion, it seems reasonable to expect in the near future accurate data also on target polarization

transfer with the advent of facilities with continuous wave electron beams.

In general, the purpose of the study of spin observables is to exploit the enhanced sensitivity of such non-averaged observables to small components of the transition amplitude in order to test specific points of the theory (potential models, meson exchange or quark effects,...) or to obtain information about badly known quantities. A very interesting example is G_E^n which could be determined via neutron polarization experiments in quasi free scattering. We have recalled above the sideways component of the beam polarization transfer. As we shall see, the longitudinal component of the target polarization transfer in scattering off vector polarized deuterons is as sensitive as $P_{0x}^h(n)$ to G_E^n models.

A problem not yet settled in the theory of the EM nuclear interactions is the choice of the EM form factors in the nucleon charge and current operators, essentially the choice of the Dirac F_1 or the Sachs G_E form factor as nucleon charge density. This uncertainty worries either IA operators and MEC operators. For long time and by many authors it has been argued that the same form factor should appear in IA current and in the (longitudinal part of) meson exchange currents so as to satisfy the continuity equation. This constraint has been questioned by Gross and Riska [29] (see also Refs.[30-31]) on the basis that the EM form factor is naturally different in different EM processes. For example, the pionic current is governed by the pion form factor F_π and the contact current by the axial form factor F_A . They have shown that it is possible to fulfill gauge invariance while remaining free to use arbitrary EM form factor in IA and MEC operators. This is achieved by adding to the standard EM vertices appropriate terms which make them able to satisfy the Ward - Takahashi identity, which is the gauge invariance constraint at the operator level. Note that these additional terms are purely longitudinal and thus vanish for on-shell particles. Moreover, they do not contribute to transition amplitudes because the electron current is conserved. These prescriptions for the off-shell extension of the EM vertices (as well as for the inclusion of the hadronic FF in the meson propagators) have been criticized [32] for their arbitrariness and also because of their limited validity. Indeed, the off-mass-shell form factors cannot depend only on Q^2 as assumed in Ref.[29]. Clearly, this proposal has given a theoretical justification also to calculations which go with G_E in IA and with F_1 in MEC. From the numerical point of view, this mixed case exhausts the possibilities for pion MEC because F_π and F_A are rather close to the isovector part of F_1 in the intermediate Q^2 range.

Some light on the problem is shed from considering the Dirac and Sachs relativistic forms of the γNN -vertex. In fact, in the non relativistic reduction one obtains in the two cases corrections of order $(v/c)^2$ which to a large extent compensate for the sizeable

differences in IA operators. For example, the Dirac (ρ^D) and the Sachs (ρ^S) charge density operators in the two-component spinor space read

$$\begin{aligned}\rho^D &= F_1 - (F_1 + 2F_2) \frac{q^2 + i\boldsymbol{\sigma} \times \mathbf{P} \cdot \mathbf{q}}{8M^2} \quad , \\ \rho^S &= G_E - G_E \frac{q^2}{8M^2} - i(2G_M - G_E) \frac{\boldsymbol{\sigma} \times \mathbf{P} \cdot \mathbf{q}}{8M^2} \quad ,\end{aligned}\tag{1}$$

for energy transfer q_0 negligible with respect to three-momentum transfer q . In this kinematic conditions the relation between G_E and F_1 becomes $G_E = F_1 - F_2 q^2 / 4M^2$ and then ρ^D and ρ^S differ only by terms of fourth order. Clearly, the F_2 part of the Darwin-Foldy term in ρ^D , a substantial part of RC in the Dirac parametrization, is already included in IA using the Sachs parametrization.

Consequently, calculations including RC should be little dependent on the choice of G_E or F_1 for processes dominated by the nucleonic degrees of freedom as the deuteron electrodisintegration in the quasi-elastic region. This has been verified by explicit evaluation of selected examples of structure functions and spin observables in Ref.[14] and by Arenhövel et al. in their systematic study of the form factors in the inclusive process [33] and of the structure functions in the exclusive process [34] (see also Ref.[35] where in addition the frame dependence of the calculations has been considered). In particular, since MEC play only a minor role at the quasi-elastic peak, the possible choices (F_1 , G_E , F_π , F_A) of the EM form factor in MEC do not give results significantly different. In conclusion, we shall mainly report results of calculations with the Dirac current and just comment on the differences with calculations with the Sachs form because, even if not definite, there are arguments in favour of the Dirac form of the γNN -vertex. It follows by minimal substitution from the Dirac Hamiltonian and at least for structureless nucleons, it fulfills the Ward-Takahashi identity. Also, it is consistent with the use of the Dirac nucleon propagator in the evaluation of MEC and of pionic RC. Moreover, it does not seem possible to formulate a chiral invariant effective Lagrangian predicting the Sachs form of the γNN -vertex [4].

Finally, we shall not deal with the potential model dependence of the polarizations. All the calculations are done with the Paris potential [36]. The reason for this is that the dependence on potential models is not substantial in the quasielastic region [22].

In Sect.2, we shall derive general formulas for cross section and nucleon polarization when both beam and target are polarized. Our results for all the components of nucleon polarization in the coplanar kinematics will be presented in Sect.3. Finally, we shall summarize our work in Sect.4.

II. FORMALISM

In this paper we shall concentrate on the outgoing nucleon polarization in exclusive deuteron electrodisintegration with both beam and target polarized. We shall restrict our considerations to the case where the deuteron target is polarized with axial symmetry with respect to an orientation axis \mathbf{d} (characterized by the angles (θ_d, ϕ_d) with respect to the virtual photon momentum \mathbf{q}) so that its state of orientation is defined by two parameters, the vector P_1^d and tensor P_2^d parameters. Note that this corresponds to the present experimental situation for polarized deuteron target with the quantization axis defined by an external magnetic field. We refer to the paper by Dmitrasinovic and Gross [37] for the general case of deuterons anyway polarized.

Arenhövel et al. [22] have already given general formulas for cross section and, with further limitation of unpolarized target, for outgoing nucleon polarization. In the following we shall use the notations of Ref.[22] apart from some minor modifications.

The polarization of the nucleon detected in coincidence with the electron is expressed in terms of the T -transition matrices and the virtual photon and deuteron density matrices by

$$\left(\frac{d\sigma}{dE_{e'} d\Omega_{e'} d\Omega_N^{cm}} \right) \mathbf{P} = \frac{\sigma_M}{M_d} \sum_{\mu\mu'} v_{\mu\mu'} \times \sum_{sm_s s' m'_s} \sum_{m_d m'_d} T_{sm_s \mu m_d(\mathbf{q})} \rho_{m_d m'_d}^d T_{s' m'_s \mu' m'_d(\mathbf{q})}^* \langle s' m'_s | \boldsymbol{\sigma} | sm_s \rangle \quad , \quad (2)$$

where $\boldsymbol{\sigma}$ is the spin operator of the nucleon. For convenience we have factorized out in Eq.(2) the Mott cross section σ_M and the deuteron mass M_d . The transition matrices $T_{sm_s \mu m_d(\mathbf{q})}$ in Eq. (2) are so denoted to recall that m_d is the deuteron spin projection on the momentum transfer \mathbf{q} . The final state is characterized by spin s and its projection m_s on the relative momentum \mathbf{p}_{cm} . We assume, as usual, relativistic electron energies so that the electron beam may only have longitudinal polarization of degree h . Therefore, $v_{\mu\mu'}$ consists of two terms, $v_{\mu\mu'} = v_{\mu\mu'}^0 + h v_{\mu\mu'}^h$, which correspond to unpolarized and polarized electrons, respectively. They are kinematical functions depending only on the electron scattering variables. Because of the symmetry relations

$$\begin{aligned} v_{\mu\mu'} &= v_{\mu'\mu} \quad , \\ v_{-\mu-\mu'}^0 &= (-)^{\mu+\mu'} v_{\mu\mu'}^0 \quad , \\ v_{-\mu-\mu'}^h &= (-)^{\mu+\mu'+1} v_{\mu\mu'}^h \quad , \end{aligned} \quad (3)$$

all the possible components can be simply derived from

$$\begin{aligned}
v_L^0 &= \left(\frac{q_{lab}}{q_{cm}} \right)^2 \xi^2 , \\
v_T^0 &= \eta + \frac{1}{2}\xi , \\
v_{TL}^0 &= \frac{1}{\sqrt{2}} \left(\frac{q_{lab}}{q_{cm}} \right) \xi \sqrt{\eta + \xi} , \\
v_{TT}^0 &= -\frac{1}{2}\xi , \\
v_T^h &= \sqrt{\eta(\eta + \xi)} , \\
v_{TL}^h &= \frac{1}{\sqrt{2}} \left(\frac{q_{lab}}{q_{cm}} \right) \xi \sqrt{\eta} ,
\end{aligned} \tag{4}$$

where the indices L , T , TL and TT correspond to $(\mu\mu') = (00)$, (11) , (10) and $(1-1)$; q_{lab} and q_{cm} are the moduli of \mathbf{q} in the laboratory (lab) and final n-p center-of-mass (cm) frames; $\xi = Q^2/q_{lab}^2$, $\eta = \tan^2(\vartheta_{e'}/2)$, $Q^2 = \mathbf{q}^2 - q_0^2$ being the four-momentum transfer squared and $\vartheta_{e'}$ the electron lab scattering angle.

Note that the definitions (4) of the v 's include the appropriate factors of (q_{lab}/q_{cm}) which are necessary because we calculate the nuclear matrix elements in the cm frame.

The dependence of the T -transition matrices on the angle ϕ between the reaction plane and the scattering plane can be separated out, defining the reduced t -matrices

$$T_{sm_s\mu m_d}(\mathbf{q}) = e^{i(\mu+m_d)\phi} t_{sm_s\mu m_d}(\mathbf{q}) . \tag{5}$$

Clearly, in order to evaluate the reduced t -matrix elements one has to refer both initial and final states to a common quantization axis. To this end we have to express the final state with respect to \mathbf{q} or the deuteron state with respect to \mathbf{p}_{cm} . In the second case we must rotate the deuteron state through the angles $(0, -\vartheta_{cm}, 0)$, ϑ_{cm} being the polar angle of \mathbf{p}_{cm} with respect to \mathbf{q} . Hence, the reduced transition matrices must be transformed according to

$$t_{sm_s\mu m_d}(\mathbf{q}) = \sum_{m'_d} t_{sm_s\mu m'_d}(\mathbf{p}_{cm}) d_{m'_d m_d}^1(-\vartheta_{cm}) . \tag{6}$$

We shall neglect in the following the specification of the quantization axis when not necessary. The t -matrix elements are expressed by means of the charge and current matrix elements as:

$$\begin{aligned}
t_{sm_s\mu m_d} &= - \sqrt{\frac{p_{cm} E_{cm}^p E_{cm}^d}{16\pi^3}} e^{-i(\mu+m_d)\phi} \\
&\times \langle sm_s | \delta_{\mu 0} \rho(\mathbf{q}) + \delta_{|\mu|1} \mathbf{e}_\mu \cdot \mathbf{j}(\mathbf{q}) | m_d \rangle ,
\end{aligned} \tag{7}$$

having denoted by \mathbf{e}_μ the photon polarization vector of helicity μ . Owing to the factorization of σ_M/M_d in Eq. (2) the reduced t -matrix is dimensionless as that introduced in Ref.[12]. The only difference is in the kinematic factors which maintain in Eq.(7) the relativistic expression. As for the final np state $|sm_s\rangle$ it is normalized so that it becomes

$$|sm_s\rangle = e^{i\mathbf{p}_{cm}\cdot\mathbf{r}}\chi_{sm_s} \quad (8)$$

in plane wave (PW) approximation.

As said above, the deuteron density matrix is diagonal with respect to the symmetry axis \mathbf{d} . Thus, its expression in terms of the orientation parameters $P_I^d(P_0^d = 1)$ is

$$\rho_{m_d m'_d}^d = \sum_{I=0}^2 \sum_{M'=-I}^{+I} \sqrt{\frac{\hat{I}}{3}} (-1)^{1-m_d} \begin{pmatrix} 1 & 1 & I \\ m_d & -m'_d & -M \end{pmatrix} e^{-iM\phi_d} d_{M0}^I(\theta_d) P_I^d \quad , \quad (9)$$

($\hat{I} = 2I + 1$) in the frame with \mathbf{q} as quantization axis taking account of the rotation which transforms P_I^d from the \mathbf{d} -frame to the \mathbf{q} -frame.

Note that the deuteron density matrix is not affected by the Lorentz transformation along \mathbf{q} , \mathbf{q} being the quantization axis of the deuteron state.

Using expressions (5)-(9) we can write the polarization of the detected nucleon in the form:

$$\left(\frac{d\sigma}{dE_e d\Omega_e d\Omega_N^m} \right) \mathbf{P} = \frac{\sigma_M}{M_d} \sum_{I=0}^2 P_I^d [\mathbf{p}_I + h\mathbf{p}_I^h] \quad . \quad (10)$$

If we define the quantity

$$\alpha_{\mu\mu'}^{IM(\pm)} = [(\mu - \mu')\phi + M(\phi - \phi_d)] \pm (I - 1)\frac{\pi}{2} \quad , \quad (11)$$

the cartesian components of the polarization vectors are given by

$$\begin{aligned} p_I^{x/z} &= \sum_{M=-I}^{+I} d_{M0}^I(\theta_d) \sum_{\mu=0}^1 \sum_{\mu'=-\mu}^{+\mu} \cos \alpha_{\mu\mu'}^{IM(-)} v_{\mu\mu'}^0 g_{\mu\mu'IM}^{x/z} \quad , \\ p_I^y &= \sum_{M=-I}^{+I} d_{M0}^I(\theta_d) \sum_{\mu=0}^1 \sum_{\mu'=-\mu}^{+\mu} \sin \alpha_{\mu\mu'}^{IM(+)} v_{\mu\mu'}^0 g_{\mu\mu'IM}^y \quad , \\ p_I^{hx/z} &= \sum_{M=-I}^{+I} d_{M0}^I(\theta_d) \sum_{\mu'=0}^1 \sin \alpha_{1\mu'}^{IM(+)} v_{1\mu'}^h g_{1\mu'IM}^{hx/z} \quad , \\ p_I^{hy} &= \sum_{M=-I}^{+I} d_{M0}^I(\theta_d) \sum_{\mu'=0}^1 \cos \alpha_{1\mu'}^{IM(-)} v_{1\mu'}^h g_{1\mu'IM}^{hy} \quad , \end{aligned} \quad (12)$$

with respect to the right-handed frame which, according to the Madison convention, has the z-axis along the nucleon momentum \mathbf{p}_{cm} and the y-axis along $\mathbf{q} \times \mathbf{p}_{cm}$.

The polarization structure functions (PSF) $g_{\mu\mu'IM}^{x/y/z}$ and $g_{\mu\mu'IM}^{hx/y/z}$ are defined by:

$$\begin{aligned}
g_{\mu\mu'IM}^x &= -\sqrt{2} \left(\frac{1 + \delta_{\mu+\mu',1}}{1 + \delta_{\mu,0}} \right) \text{Im} \left(i^I [p_{\mu\mu'IM}^1 - p_{\mu\mu'IM}^{-1}] \right) \quad , \\
g_{\mu\mu'IM}^y &= \sqrt{2} \left(\frac{1 + \delta_{\mu+\mu',1}}{1 + \delta_{\mu,0}} \right) \text{Im} \left(i^{-I} [p_{\mu\mu'IM}^1 + p_{\mu\mu'IM}^{-1}] \right) \quad , \\
g_{\mu\mu'IM}^z &= 2 \left(\frac{1 + \delta_{\mu+\mu',1}}{1 + \delta_{\mu,0}} \right) \text{Im} \left(i^I p_{\mu\mu'IM}^0 \right) \quad , \\
g_{\mu\mu'IM}^{hx} &= \sqrt{2} \left(\frac{1 + \delta_{\mu+\mu',1}}{1 + \delta_{\mu,0}} \right) \text{Re} \left(i^{-I} [p_{\mu\mu'IM}^1 - p_{\mu\mu'IM}^{-1}] \right) \quad , \\
g_{\mu\mu'IM}^{hy} &= \sqrt{2} \left(\frac{1 + \delta_{\mu+\mu',1}}{1 + \delta_{\mu,0}} \right) \text{Re} \left(i^I [p_{\mu\mu'IM}^1 + p_{\mu\mu'IM}^{-1}] \right) \quad , \\
g_{\mu\mu'IM}^{hz} &= -2 \left(\frac{1 + \delta_{\mu+\mu',1}}{1 + \delta_{\mu,0}} \right) \text{Re} \left(i^{-I} p_{\mu\mu'IM}^0 \right) \quad ,
\end{aligned} \tag{13}$$

in terms of the p -functions

$$\begin{aligned}
p_{\mu\mu'IM}^\lambda &= \sqrt{\frac{\hat{I}}{3}} \sum_{M'=-I}^{+I} \sum_{m_d m'_d} (-1)^{1-m_d} d_{M'M}^I (-\vartheta_{cm}) \begin{pmatrix} 1 & 1 & I \\ m_d & -m'_d & -M' \end{pmatrix} \\
&\times \sum_{sm_s s' m'_s} t_{sm_s \mu m_d} t_{s' m'_s \mu' m'_d}^* \langle s' m'_s | \sigma_\lambda | sm_s \rangle \quad ,
\end{aligned} \tag{14}$$

where σ_λ is the $\lambda = 0, \pm 1$ spherical component of the spin operator.

Clearly Eq.(12) exhibits polarization components in a very symmetric and compact form which however suffers from a minor and a major drawback. The minor one is that the dependence on the azimuthal angle ϕ is not explicit. The major one is that these components are written in terms of more PSF than necessary. In fact, the PSF's (13) are not all independent and even not all different from zero because of the symmetry relations fulfilled by the p -functions:

$$\begin{aligned}
p_{\mu\mu'IM}^{\lambda*} &= (-1)^{\lambda+M} p_{\mu'\mu I-M}^{-\lambda} \quad , \\
p_{-\mu-\mu'IM}^\lambda &= (-1)^{1+I+M+\lambda+\mu+\mu'} p_{\mu\mu' I-M}^{-\lambda} \quad .
\end{aligned} \tag{15}$$

The first relation in Eq.(15) comes directly from definition (14) while the second one needs in addition the symmetry relation induced on the t -matrix elements by parity conservation

$$t_{s-m_s-\mu-m_d} = (-1)^{1+s+m_s+\mu+m_d} t_{sm_s \mu m_d} \quad . \tag{16}$$

Note that equalities (15) have already been used to reduce the summation over μ and μ' to the four cases denoted L, T, TL, TT . Furthermore, it is straightforward to deduce from Eq.(15) that the non-interference PSF with $M \leq 0$ differ at most for the sign from those PSF with $M \geq 0$. Explicitly we have :

$$\begin{aligned} g_{\mu\mu, I-M}^{x/z} &= (-1)^{I-M+1} g_{\mu\mu, IM}^{x/z} , \\ g_{\mu\mu, I-M}^y &= (-1)^{I-M} g_{\mu\mu, IM}^y , \\ g_{\mu\mu, I-M}^{h \ x/z} &= (-1)^{I-M} g_{\mu\mu, IM}^{h \ x/z} , \\ g_{\mu\mu, I-M}^{h \ y} &= (-1)^{I-M+1} g_{\mu\mu, IM}^{h \ y} , \end{aligned} \quad (17)$$

for $\mu=0,1$. Therefore, the independent PSF are 121 (40 with $M = 0$ and 81 with $M \neq 0$). The use of relations (16) and (17) allows to partially reduce the summation over M in some but not in all the polarization components, with the consequence that the high symmetry of expression (12) is lost. For this reason we prefer not to further handle these expressions.

For completeness, we also report the cross section in a similarly compact form, namely

$$\frac{d\sigma}{dE_{e'} d\Omega_{e'} d\Omega_N^{cm}} = \frac{\sigma_M}{M_d} \sum_{I=0}^2 P_I^d (F_I + h F_I^h) \quad . \quad (18)$$

The functions F_I, F_I^h are given by

$$\begin{aligned} F_I &= \sum_{M=-I}^{+I} d_{M0}^I(\theta_d) \sum_{\mu=0}^1 \sum_{\mu'=-\mu}^{+\mu} \sin \alpha_{\mu\mu'}^{IM(+)} v_{\mu\mu'}^0 f_{\mu\mu' IM} , \\ F_I^h &= \sum_{M=-I}^{+I} d_{M0}^I(\theta_d) \sum_{\mu'=0}^1 \cos \alpha_{1\mu'}^{IM(-)} v_{1\mu'}^h f_{1\mu' IM}^h , \end{aligned} \quad (19)$$

in terms of the structure functions (SF)

$$\begin{aligned} f_{\mu\mu' IM} &= -2 \left(\frac{1 + \delta_{\mu+\mu',1}}{1 + \delta_{\mu,0}} \right) \text{Re}(i^{-I} w_{\mu\mu' IM}) , \\ f_{\mu\mu' IM}^h &= 2 \left(\frac{1 + \delta_{\mu+\mu',1}}{1 + \delta_{\mu,0}} \right) \text{Im}(i^I w_{\mu\mu' IM}) , \end{aligned} \quad (20)$$

where

$$\begin{aligned} w_{\mu\mu' IM} &= \sqrt{\frac{\hat{I}}{3}} \sum_{M'=-I}^{+I} \sum_{m_d m'_d} (-1)^{1-m_d} d_{M'M}^I(-\vartheta_{cm}) \begin{pmatrix} 1 & 1 & I \\ m_d & -m'_d & -M' \end{pmatrix} \\ &\times \sum_{sm_s} t_{sm_s \mu m_d} t_{sm_s \mu' m'_d}^* . \end{aligned} \quad (21)$$

Here again there are symmetry relations among the auxiliary functions $w_{\mu\mu'IM}$:

$$\begin{aligned} w_{\mu\mu'IM}^* &= (-1)^M w_{\mu'\mu I-M} \quad , \\ w_{-\mu-\mu'IM} &= (-1)^{I-M+\mu+\mu'} w_{\mu\mu' I-M} \quad , \end{aligned} \quad (22)$$

which have already been used to write Eq.(19) in terms of only $\mu = 0, 1; -\mu \leq \mu' \leq \mu$.

Analogously to the symmetry relations (15) for the p -functions, the first equality (22) is an immediate consequence of definition (21), while the second one follows from the parity conservation relation (16). Owing to Eq.(22)

$$\begin{aligned} f_{\mu\mu, I-M} &= (-1)^{I-M} f_{\mu\mu, IM} \quad , \\ f_{\mu\mu, I-M}^h &= (-1)^{I-M+1} f_{\mu\mu, IM}^h \quad , \end{aligned} \quad (23)$$

and therefore the 54 SF appearing in Eq.(20) reduce to 41 independent SF (14 with $M = 0$ and 27 with $M \neq 0$). The explicit expression of the cross section in terms of the independent SF can be found in Ref.[34].

The usual definition of the polarization variables is

$$\left(\frac{d\sigma}{dE_e d\Omega_{e'} d\Omega_N^{cm}} \right) \mathbf{P} = \left(\frac{d\sigma}{dE_e d\Omega_{e'} d\Omega_N^{cm}} \right)_0 \sum_{I=0}^2 P_I^d [\mathbf{P}_I + h \mathbf{P}_I^h] \quad , \quad (24)$$

where $(d\sigma/dE_e d\Omega_{e'} d\Omega_N^{cm})_0 = (\sigma_M/M_d)F_0$ is the cross section with unpolarized beam and target. Therefore, we have the following relation between the \mathbf{p}' s and \mathbf{P}' s

$$\begin{aligned} F_0 \mathbf{P}_I &= \mathbf{p}_I \quad , \\ F_0 \mathbf{P}_I^h &= \mathbf{p}_I^h \quad . \end{aligned} \quad (25)$$

Since the outgoing nucleon polarization depends on six independent vectors, one has in general, as many as 18 different components to study. They halve in the in-plane kinematics if the deuteron symmetry axis \mathbf{d} lies in the reaction plane, case to which we confine ourselves with the choice $\theta_d = 90^\circ$, $\phi_d = 0^\circ$. In this case only the component of \mathbf{P}_0 , \mathbf{P}_2 , and \mathbf{P}_1^h normal to the reaction plane and the two components of \mathbf{P}_0^h , \mathbf{P}_1 , and \mathbf{P}_2^h lying in the reaction plane are non vanishing because of parity conservation. In fact, the electron helicity h and the vector orientation parameter P_1^d change sign while the tensor orientation parameter P_2^d remains unchanged under space inversion.

In principle the separation of the different polarization vectors in Eq.(24) from the unavoidable \mathbf{P}_0 is a simple task, dictated as it is by the range of values covered by h , P_1^d , and P_2^d . We recall that $h = \pm 1$ for relativistic electrons and that $\sqrt{\frac{2}{3}}P_1^d = n_1 - n_{-1}$ and $\sqrt{2}P_2^d = 1 - 3n_0$, where n_m is the fraction of the deuterons with spin projection m

on the symmetry axis \mathbf{d} . Therefore, $\sqrt{\frac{2}{3}}P_1^d$ varies in the range $(-1, 1)$ and $\sqrt{2}P_2^d$ in the range $(-2, 1)$. Clearly, to determine \mathbf{P}_0^h one has to perform two measurements of \mathbf{P} with unpolarized target and with opposite values of the electron helicity and to subtract the results. Obviously such a double experiment is necessary in the general case but not in the in-plane kinematics where \mathbf{P}_0 and \mathbf{P}_0^h are orthogonal, as noted above. Similarly, \mathbf{P}_1 comes from the difference of \mathbf{P} measured with unpolarized beam and vector polarized deuteron target with $\pm P_1^d$. To obtain \mathbf{P}_1^h one needs four experiments with the four combinations of $\pm h$ and $\pm P_1^d$ and the appropriate sums and differences of the resulting polarizations. The same procedure allows to disentangle \mathbf{P}_2 and \mathbf{P}_2^h because also $\sqrt{2}P_2^d$ can assume opposite values in the range $(-1, 1)$.

Much more complicated is the problem of disentangling the different PSF defining the polarization vectors (see Eq.(12)). It usually requires several measurements in different kinematic conditions, at different values of polarization parameters and also exploiting the dependence of the polarization components on the polar and azimuthal angles of \mathbf{d} . The similar problem concerning the various SF of the cross section has been discussed in Ref.[34]. Since the determination of PSF is beyond the present experimental feasibility we concentrate our discussion directly on the polarization vectors.

Finally, we address the question of the transformation to the *lab* frame of cross section and nucleon polarization. In order to have the differential cross section fully in the *lab* frame it is sufficient to multiply Eq.(18) by the Jacobian

$$\frac{\partial \Omega_p^{cm}}{\partial \Omega_N^{lab}} = \frac{1}{\gamma} \frac{\left[\gamma^2 (\beta/\beta^{cm} + \cos \vartheta^{cm})^2 + \sin^2 \vartheta^{cm} \right]^{3/2}}{1 + (\beta/\beta^{cm}) \cos \vartheta^{cm}}, \quad (26)$$

because the differentials of the electron energy and solid angle in Eq.(18) are already *lab* variables. In Eq.(26) β^{cm} is the *cm* nucleon velocity, β and γ are the boost parameters.

The transformation of the nucleon polarization is less trivial because the nucleon momentum is in general not colinear with the boost. Therefore, the polarization vector undergoes a Wigner rotation [38] about the axis $\mathbf{q} \times \mathbf{p}_{cm}$ through an angle ψ which can be put in the form

$$\tan \psi = \frac{\sin \vartheta^{cm}}{\gamma^{cm} (\beta^{cm}/\beta + \cos \vartheta^{cm})}. \quad (27)$$

Hence, the boost leaves unchanged the *y* component of the polarization but mixes the *x* and *z* components. Explicitly, one has:

$$\begin{aligned} P_s^{lab} &= P_x \cos \psi + P_z \sin \psi, \\ P_l^{lab} &= -P_x \sin \psi + P_z \cos \psi, \end{aligned} \quad (28)$$

where P_l^{lab} and P_s^{lab} are the longitudinal and sideways *lab* components in the $\phi = 0$ half-plane. Note that there is no mixing for forward and backward emitted nucleons.

III. RESULTS

We shall consider the observables in three theoretical approximations, the IA theory which corresponds to the standard NR calculations with nucleonic only contributions to the reaction amplitude and final state interaction (FSI) included; the IA+MEC+IC theory which takes into account the mesonic exchange currents and the Δ -excitation current of pionic range, the last one in the static approximation for the Δ propagator; the full theory which also includes the nucleonic relativistic corrections to the EM operators as well as to the wave functions and the pionic RC to the charge density in PV coupling theory. In actual calculations we use the Paris potential, as said in the Introduction, and for the nucleon EM form factor we use the dipole fit and the Galster model [39] of G_E^n (with $p=5.6$). The sensitivity to G_E^n models is investigated by also considering the same dipole fit and $G_E^n = 0$.

For more details about meson exchange and relativistic corrections we refer to Ref.[12] where the technical problem caused by the low convergence of the multipole expansion of the transition amplitude for the nucleonic contributions is also discussed. The short range of the mesonic operators ensures a good convergence of the mesonic parts of the t -matrix. As a result, we explicitly evaluate Coulomb, electric, and magnetic multipoles up to order $L = 6$ with FSI while all the higher multipoles are calculated in PW approximation. This is obtained evaluating the PW amplitude in closed form, i.e. without recourse to the multipole expansion, as first shown by Renard et al. [40] and widely discussed by Fabian and Arenhövel [41]. In order to avoid double counting one has to subtract the first multipoles calculated in PW approximation. To be fully consistent in such a procedure, we have chosen the Foldy gauge [42] (see also Friar and Fallieros [43]) in all calculations (closed form in PW approximation and multipole expansion without and with FSI). We would like to note that the use of the Siegert form of the transverse electric multipoles allows us to also account for the major part of the pionic RC through the knowledge of the pionic charge density.

For simplicity, we do not report results obtained in PW approximation for the final states. We recall that in such an approximation the reduced transition matrix is real in compliance with Watson's final state theorem. As a consequence, several components of the polarization vectors and notably \mathbf{P}_0 , vanish when FSI are switched off because they depend on the imaginary part of appropriate p -functions (see Eqs. (13-14)). The inadequacy of this approximation in describing the angular shape of the other components, which are given by the real part of PSF, has been already demonstrated in Ref.[22-24]. A noticeable exception which is worth mentioning, is the forward direction. Here, as shown by Arenhövel et al. [22], these polarization components are essentially

given by PWIA.

As kinematical region we have chosen the quasi-free scattering (defined by the Bjorken variable $x = 1$) at $Q^2 = 12 fm^{-2}$ and then for relative n-p energy of about 120 MeV, a condition which maximizes the momentum transfer involved and therefore the importance of RC while remaining below the pion production threshold where the description of the NN interaction in terms of the realistic NN potentials is valid. To completely fix the kinematics we assume the electron scattering angle $\vartheta_e = 60^\circ$. The corresponding electron beam energy is of 820 MeV and the scattered electron energy of 570 MeV, namely a situation well within the possibilities of the existing accelerators. For the polarization vector we take a reference frame slightly different from that in Sect. 2. More precisely, we leave unchanged the z-axis in the direction of the nucleon momentum, but we take the y-axis parallel to $\mathbf{k} \times \mathbf{k}'$, \mathbf{k} and \mathbf{k}' being the initial and final electron momenta. Had we followed the Madison convention of Sect. 2. with the y-axis directed along $\mathbf{q} \times \mathbf{p}_{cm}$, the y-components would have shown an unpleasant change of sign for nucleon emission to the right or left of \mathbf{q} . Polarizations are drawn as a function of the nucleon cm polar angle measured from the virtual photon direction. Therefore, $\vartheta_{cm} = 0^\circ$ is the strict quasi elastic peak. The sign of ϑ_{cm} has been assumed positive in the half-plane $\phi = 0^\circ$ and negative in the half-plane $\phi = 180^\circ$.

We start considering the nucleon polarization in the unpolarized case. As already said, only the component of \mathbf{P}_0 normal to the reaction plane survives in the in-plane kinematics because of parity conservation. The first remark worth doing is that the nucleon polarization takes on appreciable values which make it accessible for measurements.

The proton polarization $P_{0y}(p)$ reported in Fig. 1b shows a remarkable sensitivity to both MEC and RC for all angles except a narrow forward region. In particular these effects add coherently leading to a reduction by a factor 2 of the maximum centered at $\vartheta_{cm} = -50^\circ$ and to the disappearance of the peak at $\vartheta_{cm} = 70^\circ$. Also the pronounced maximum at $\vartheta_{cm} = 150^\circ$ suffers from a 15% lowering. At higher angles in the half-plane $\phi = 180^\circ$ MEC and RC interfere destructively resulting in a small and oscillating polarization.

The neutron polarization $P_{0y}(n)$ is drawn in Fig. 1a. It is characterized in IA theory by a sharp maximum followed by a bump and a deep minimum in the left half-plane and by a large maximum with a peak at $\vartheta_{cm} = 130^\circ$ in the right half-plane. While the minimum is slightly modified by MEC and RC, the shape of the two maxima is heavily affected. In particular, the bump at $\vartheta_{cm} = -110^\circ$ and the peak at $\vartheta_{cm} = 130^\circ$ reduce to two shoulders owing to the coherent MEC and RC effect.

These results are for calculations with F_1 . If G_E is used instead, the angular behaviour of the nucleon polarization remains essentially the same with its rich structure of maxima and minima. However, the absolute values show sizeable differences both in the IA predictions which are smaller on the average and in the IA+MEC+IC theory. With inclusion of RC one gets full theory curves very similar to those in Fig. 1.

If the electron beam is polarized and the target unpolarized, the nucleon polarization vector acquires a new contribution, the beam polarization transfer \mathbf{P}_0^h , which is parallel to the reaction plane in the in-plane kinematics. The higher values of the nucleon polarization are reached by the longitudinal component P_{0z}^h in the forward direction as shown in Fig. 2. Two other maxima appear at backward angles in the two half planes. Both protons and neutrons present a positive polarization transfer, i.e. they are emitted with spin aligned with the direction of motion. The IA theory fix the shape of $P_{0z}^h(n)$ in the whole angular range while the IA+MEC+IC results are close to the full theory values. Indeed, the relativistic effect is visible in the shift towards the right of the large forward peak and in backward direction. Instead, RC play a very important role in defining the longitudinal polarization of the proton either in the forward peak (with a 30% increase at $\vartheta_{cm} = 0^\circ$) and in the half-plane $\phi = 180^\circ$. As for the calculations with G_E we would like to remark that the relativistic effect is weaker but still noticeable in $P_{0z}^h(p)$, particularly in the backward peaks and at 0° where it amounts to a 10% increase.

The sideways component P_{0x}^h of \mathbf{P}_0^h shows in general less pronounced but still measurable values with the noticeable exception of the deep minimum at $\vartheta_{cm} = 30^\circ$ for emitted protons, Fig. 3b. Such minimum is largely due to the nucleonic NR contributions. Mesonic and relativistic corrections are effective at higher angles where however, $P_{0x}^h(p)$ assumes rather small values. Instead, RC considerably modify P_{0x}^h for emitted neutrons, Fig. 3a, in the forward region with an upward shift of the polarization which even changes sign in parallel kinematics. We recall that a measurement of $P_{0x}^h(n)$ has been suggested by Arnold, Carlson and Gross [18] for an experimental determination of G_E^n because of its sensitivity to the different models of G_E^n . The additional curve in Fig. 3a, corresponding to calculation with $G_E^n = 0$, sets the scale of such sensitivity. The dependence of $P_{0x}^h(n)$ on potential models and on exchange effects has been shown to be very small at $\vartheta_{cm} = 0^\circ$ [21,22]. We can add that it is also little sensitive to the choice of the Dirac or Sachs form of the nucleon EM form factors. Indeed, the full theory results with G_E [14] are almost identical to those in Fig. 3. We may also note that here the RC effect is generally negligible and that it almost completely vanishes in parallel kinematics in $P_{0x}^h(n)$. Then the IA+MEC predictions for $P_{0x}^h(n)$ of Ref.[21-22] obtained with the Sachs parametrization almost coincide with our full theory results

around $\vartheta_{cm} = 0^\circ$. In order to compare our results with those of Ref.[21-22] one has to recall that Arenhövel et al. report neutron polarization as function of the proton polar angle and for neutron emitted in the $\phi = 180^\circ$ half-plane. Moreover, there is a difference of sign due to our convention for the x component.

In the case of unpolarized beam and polarized target, the nucleon polarization takes contribution from two new vectors, (in short, vector and tensor target polarization transfer) \mathbf{P}_1 and \mathbf{P}_2 which correspond to the possible vector and tensor deuteron polarization. Such polarization vectors have not been studied till now in the literature, probably because their measurement seemed quite unlikely. As discussed in the Introduction, this is no longer the case because of the continuous improvements in nucleon polarimeters and in the production of polarized internal target which, together with the availability of continuous wave electron beams, make it possible to foresee experiments of these observables in the near future. For this reason we report their non-vanishing components (recall that \mathbf{P}_1 lies in the reaction plane while \mathbf{P}_2 is normal to that plane in the coplanar kinematics) for each possible detected nucleon. The more interesting components of the target polarization transfers are the longitudinal component for outgoing neutron and the sideways component for outgoing proton in the case of vector polarized deuterons. In particular $P_{1z}(n)$, Fig. 4a, seems as promising as $P_{0x}^h(n)$ for an experimental determination of G_E^n even if the measurement of a longitudinal polarization presents an additional difficulty. Indeed, it must be turned into a transverse polarization to be measured with the focal plane polarimeters. In fact, $P_{1z}(n)$ shows the same sensitivity of $P_{0x}^h(n)$ to G_E^n in the forward direction as the comparison between the full theory curves with $G_E^n = 0$ and $G_E^n \neq 0$ clearly demonstrates. This is easily understood looking at the analytic expression for $P_{1z}(n)$ in PWIA and neglecting the D-wave component of the deuteron state. In this limit $P_{1z}(n)$ which depends in general on four PSF (see Eq. (12)), is determined at 0° only by the longitudinal-transverse PSF, i.e. by the interference between charge and magnetic neutron form factors. We want to note that our results for $G_E^n = 0$ do not vanish, as one could naively expect, not only because of the influence of MEC and FSI and the D-wave deuteron component but also because our calculations are with the Dirac form of nucleon current (obviously $F_1 \neq 0$ even for $G_E = 0$). Note that similarly to $P_{0x}^h(n)$, $P_{1z}(n)$ is considerably affected by RC in the forward region where the inclusion of RC causes a change of sign. The higher values of $P_{1z}(n)$ are attained at higher angles in both half-planes, where RC are still of importance. In particular, the relativistic effect completely cancels the exchange effect in the maximum at $\vartheta_{cm} = 60^\circ$, where the sensitivity to G_E^n is still considerable. Also in this observable full theory calculations with G_E give very similar results even if the relativistic effect is much less pronounced.

The results corresponding to the Sachs form factors as reported in Fig. 5, show that the large difference at $\vartheta_{cm} = 0^\circ$ between calculations with $G_E^n = 0$ and $G_E^n \neq 0$ does not depend on the nucleon current parametrization. Here, however, the $G_E^n = 0$ curve is almost vanishing, so demonstrating the small influence of FSI and MEC and of the D-deuteron state at this point. From Fig. 5 it also clearly appears that RC are less effective in calculations with G_E . A few remarks are worth doing concerning the different role played by mesonic and relativistic corrections in the two calculations. With reference to the peak centered around $\vartheta_{cm} = 0^\circ$, we may note that an almost identical final result is obtained through a large negative MEC effect balanced by a large positive RC effect in calculations with F_1 . Instead, in calculations with G_E , the higher IA values are considerably lowered by MEC and very slightly increased by RC. Also remarkable is the fact that the relativistic effect may even change sign in the two cases, as is evident in the large peak centered at $\vartheta_{cm} = -100^\circ$. Mesonic and relativistic effects are small in $P_{1z}(p)$ in both calculations with F_1 (Fig. 4b) and G_E (Fig. 5b) except in the region from $\vartheta_{cm} = -70^\circ$ to $\vartheta_{cm} = -140^\circ$ where MEC+IC contribute significantly.

The sideways proton polarization $P_{1x}(p)$, which is reported in Fig. 6b, shows considerable variations due to RC over a large range of ϑ_{cm} . In fact, a rather strong relativistic effect is present not only in the forward region but also in the two deep minima at medium angles. It is worth mentioning that even if not so marked the RC effect remains visible also in G_E calculations. Unlike $P_{1z}(n)$, $P_{1x}(n)$ (Fig. 6a) is little sensitive to G_E^n and also little affected by MEC and RC.

The normal component of the tensor target polarization transfer P_{2y} is reported in Fig. 7. For both proton and neutron it is characterized by positive and negative peaks where its magnitude reaches values close to 1. Such values diminish in the minima and increase in the maximum when RC are considered as a consequence of a general upward shift induced by such corrections. The exchange effect is instead very small except in $P_{2y}(p)$ around $\vartheta_{cm} = -100^\circ$. Also the dependence on G_E^n is quite negligible.

It still remains to be considered the case where both beam and target are polarized. The recoil nucleon polarization takes on two new contributions denoted \mathbf{P}_1^h and \mathbf{P}_2^h in Eq.(24) which can be called vector and tensor target - beam polarization transfers. Since the measurement of nucleon polarization in such conditions is certainly beyond the range of the present experimental possibilities, we shall only briefly comment on the results of our calculations. \mathbf{P}_1^h , which is normal to the reaction plane, assumes very small values in the whole angular range for both proton and neutron emission, and is little sensitive to exchange and relativistic corrections (Fig. 8). Instead, \mathbf{P}_2^h which lies in the reaction plane, can reach appreciable values in the longitudinal (Fig. 9) and

sideways (Fig. 10) components. Whereas $P_{2x}^h(n)$ is mainly determined by IA results, $P_{2x}^h(p)$ presents interesting features at backward angles, where it shows appreciable sensitivity to RC and especially to G_E^n models as revealed by the additional curve corresponding to full theory calculations with $G_E^n = 0$. This sensitivity of $P_{2x}^h(p)$ on G_E^n has the same origin as that previously found for $P_{1z}(n)$. Indeed, a PWIA calculation with the dominant S-wave deuteron component shows that $P_{2x}^h(p)$ at 180° depends linearly on the charge form factor of the neutron. Unfortunately, in the quasi elastic peak region, the protons emitted at 180° in the cm frame move in forward direction in the lab frame with very small kinetic energy. Consequently, a measurement of their polarization is extremely difficult.

Finally, P_{2z}^h has a characteristic structure with four deep minima where the exchange effect is more pronounced, with particular evidence in the neutron case. Also RC affect P_{2z}^h in the minima at $\vartheta_{cm} = -150^\circ$ and $\vartheta_{cm} = 40^\circ$. Once again the relativistic effect considerably diminishes in the Sachs scheme.

IV. CONCLUSIONS

We have considered outgoing nucleon polarization in the case of beam and target polarized, giving the general formulas for the components, valid in any kinematic conditions. In the numerical applications, we have studied all the non-vanishing components of proton and neutron polarization in the coplanar kinematics and in the quasi-elastic region assuming the deuteron symmetry axis in the reaction plane.

Our calculations have been done in the standard theory with inclusion of the most relevant relativistic corrections. We have found that such corrections give substantial modifications in almost every polarization component. This does not come out as a surprise because we have used the Dirac form of the nucleon current in our calculations. As previous works on coincidence disintegration indicate, the size of the relativistic effect is normally larger with this form than with the Sachs form but the full theory results are very close to each other. Having verified that this is indeed the case in all the observables presented, we have chosen to report only one set of theoretical results with the exception of P_{1z} . As noted in discussing $P_{1z}(n)$, the similarity of full theory results turns out from a complicated interplay of the IA, MEC and RC contributions which are seldom individually different in size and even in sign. The fact that the predictions are little dependent on the parametrization of the nucleon current greatly reduces the ambiguity in the theoretical analysis of deuteron electrodisintegration at the quasi elastic peak.

In particular, we have studied the dependence of nucleon polarization on the neutron charge form factor, comparing results obtained with $G_E^n = 0$ and with the dipole form factor model of Galster et al. [39]. Apart from reobtaining the well known results of the sensitivity of $P_{0x}^h(n)$ on G_E^n in parallel kinematics we have picked out two new observables as sensitive as $P_{0x}^h(n)$ on G_E^n models. The first one is the longitudinal component of the vector target polarization transfer, again in forward direction. A measurement of $P_{1z}(n)$ seems to be well within the range of the experimental possibilities in the very near future in view of the recent advances in the polarized target technique and in the focal plane polarimeters. The second observable is the sideways component of tensor target - beam polarization transfer for emitted protons at $\vartheta_{cm} = 180^\circ$ which constitutes instead a prohibitively difficult experiment.

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FIGURE CAPTIONS

- Fig. 1** Outgoing neutron (a) and proton (b) polarization P_{0y} as a function of the corresponding polar angle. Dot-dashed line : IA theory; dotted line: IA + MEC + IC theory; solid line: full theory. Calculations with the Paris potential, with the Dirac form of nucleon current and with the dipole Galster model of EM form factors.
- Fig. 2** Beam polarization transfer P_{0z}^h . Notation as in Fig.1.
- Fig. 3** Beam polarization transfer P_{0x}^h . Notation as in Fig.1. The additional curve (dashed line) corresponds to $G_E^n = 0$.
- Fig. 4** Vector target polarization transfer P_{1z} . Notation as in Fig.3.
- Fig. 5** The same as in Fig.4 with the Sachs form of nuclear current.
- Fig. 6** Vector target polarization transfer P_{1x} . Notation as in Fig.1.
- Fig. 7** Tensor target polarization transfer P_{2y} . Notation as in Fig.1.
- Fig. 8** Vector target - beam polarization transfer P_{1y}^h . Notation as in Fig.1.
- Fig. 9** Tensor target - beam polarization transfer P_{2z}^h . Notation as in Fig.1.
- Fig. 10** Tensor target - beam polarization transfer P_{2x}^h . Notation as in Fig.3.